

Superconformal approach to Higgs inflation from new Kähler potential

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We introduce a new class of models of Higgs inflation using the superconformal approach to supergravity by modifying the Kähler geometry. Using such a mechanism, we construct a phenomenological functional form of a new Kähler potential. From this we construct various types of models which are characterized by a superconformal symmetry breaking parameter χ , and depending on the numerical values of χ we classify all of the proposed models into three categories. Models with minimal coupling are identified by $\chi = \pm \frac{2}{3}$ branch which are made up of shift symmetry preserving flat directions. We also propose various other models by introducing a non-minimal coupling of the inflaton field to gravity described by $\chi \neq \pm \frac{2}{3}$ branch. We employ all these proposed models to study the inflationary paradigm by estimating the major cosmological observables and confront them with recent observational data from WMAP9 and other complementary data sets. We also mention an allowed range of non-minimal couplings and the *Yukawa* type of couplings appearing in the proposed models used for cosmological parameter estimation.

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I. INTRODUCTION

Cosmological inflation has been a paradigm in which the pathological problems of the Standard Big Bang Cosmology are addressed in a sophisticated way. The inflaton generates scale-dependent nearly Gaussian spectrum of density fluctuations. Furthermore, quantum fluctuations during inflation provides a seed for the large-scale structure formation as we observe today. Inflation is governed by a flat potential which has a proper field theoretic origin [1–3]. In this context, supersymmetry or its local extension (i.e. supergravity) is the most successful candidate, which imposes certain constraints on the non-supersymmetric models of particle physics and cosmology [4–7]. A well known example of such restrictions is the fact that the supersymmetric version of the Standard Model (SM) of particle physics requires at least two Higgs superfields [8, 9]. On the other hand, the supersymmetry embedding of the Higgs model inflation requires supergravity [10–13]. Thus, it is interesting to see how supersymmetry may affect various inflationary models, where the gravity sector is minimally or non-minimally coupled to scalar fields.

An efficient proposal is to utilize the SM Higgs doublet as the inflaton via well known Higgs inflation in the context of supergravity [10–13]. In this scenario, inflation can be realized due to a large non-minimal coupling of the Higgs doublet to the Einstein gravity, instead of having a tiny Higgs quartic coupling, which is contradictory to

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the recently observed Higgs mass bound at *Large Hadron Collider* (LHC) [14]. However, it has been shown earlier by power counting formalism that the Hubble scale during inflation is proximate to the unitarity bound on the new physics scale associated with the breakdown of the semi-classical approximation in the effective field theory [15–18]. Nonetheless, a singlet field with non-minimal coupling can be a possible inflaton candidate for a small singlet self interaction motivated quartic coupling for which the Hubble scale can be smaller than the unitarity bound. Weak scale supersymmetry [8, 19] is a solution to the well posed hierarchy problem in SM and has been one of main topics in the search for new physics at the Large Hadron Collider (LHC) or any future linear collider. In the context of Minimal Supersymmetric Standard Model (MSSM) [3, 8, 19, 20], there exist two Higgs doublets and the corresponding self coupling can be expressed in terms of the electroweak gauge couplings. However, apart from the unitarity problem in the Higgs inflation, it has been shown that in MSSM, the Higgs inflation cannot be implemented without fine tuning [21] due to the instability appearing in the ratio of two Higgs VEVs. On contrary, in the context of Next-to-Minimal Supersymmetric Standard Model (NMSSM) [1, 10, 11, 22], an additional self coupling can be introduced by the superpotential term which can provide the vacuum energy required for inflation. Since this new self coupling can be made small without violating the new *LHC* bound on the Higgs mass, there is also a possibility for the Higgs inflation within the semi-classical limit.

However the supergravity theory has a dark side in the context of Higgs inflation. The main problem was rooted in the functional form of the Kähler potential which involves typical contributions proportional to quadratic combination of the superfields in the canonical version. One elegant way to overcome such problem is to search for shift symmetry [12, 23–25] protected flat directions in supergravity which can take part in inflation. The flatness of the potential is broken only by the implementation of the superconformal symmetry breaking parameter in the supergravity Kähler potential [10, 11]. Such terms are directly connected with non-minimal interactions of the inflaton field to the Einstein gravity sector. This class of non-minimal models of Kähler potential has many interesting features, which were explored in the context of the superconformal approach to supergravity. Specifically, in the context of canonical superconformal supergravity (CSS) models [10, 11] kinetic terms in the preferred frame of reference, in other words, Jordan frame are canonical and the corresponding potential is exactly same as appearing in global supersymmetry. For this purpose, in this article we propose a phenomenological model of a new Kähler potential with two singlet chiral superfields (H, S) which successfully address the problems of the supergravity inflation with non-minimal coupling (ξ_1, ξ_2). Here one singlet field plays the role of inflaton and the other one is the background which will trigger preheating [26, 27]/reheating [28–30] depending on the branching ratios of different decay channels of the inflaton. Our result can be applied directly to the Higgs inflation by satisfying D-flat constraints. In this article, our primary target is to do a thorough survey of inflationary models from Kähler potential using superconformal transformation followed by confrontation with latest observational data from *WMAP9* and other complementary datasets.

The paper is organized as follows. We first explain a general framework for $\mathcal{N}=1, \mathcal{D}=4$ Jordan frame supergravity where superconformal symmetry breaking parameters for scalar fields are suitably implemented. Then we introduce a new phenomenological model of Kähler potential with two singlet chiral superfields. Next we discuss the implication of the Higgs inflation from various types of inflationary potentials derived from the Jordan frame Kähler potential for four distinct physical branches of the symmetry breaking parameter (χ). Next imposing the constraints from *LHC* we employ these models for cosmological parameter estimation by using a numerical code *CAMB* [31]. Finally, we confront the cosmological observables with the latest *WMAP9* [32] data combining with complementary datasets in Λ CDM background.

II. SUPERCONFORMAL MECHANISM IN KÄHLER GEOMETRY

In this section we start our discussion with $\mathcal{N}=1, \mathcal{D}=4$ SUGRA action in the *Jordan frame* with generalized frame function $\Phi(z, \bar{z})$ in the Planckian unit described by [10, 11]

$$S_\Phi = \int d^4x \sqrt{-g_J} \left[R_{(4)} - 2\Lambda_{(4)} + e_{(4)}^{-1} \mathcal{L}_{SUGRA}^\Phi \right] \quad (2.1)$$

where

$$e_{(4)}^{-1} \mathcal{L}_{SUGRA}^\Phi := -\frac{\Phi(z, \bar{z})}{6} [R_{(4)} - \bar{\Psi}_\mu R^\mu] - \frac{1}{6} (\partial_\mu \Phi) (\bar{\Psi}^\alpha \gamma_\alpha \Psi^\mu) + \mathcal{L}_0 + \mathcal{L}_{\frac{1}{2}} + \mathcal{L}_1 + \mathcal{L}_m + \mathcal{L}_{mix} + \mathcal{L}_d + \mathcal{L}_{4f} - V_J. \quad (2.2)$$

In equation(2.2) the notations used are: $\Psi_\mu \Rightarrow$ gravitino field, $R^\mu \Rightarrow$ gravitino kinetic term, $\mathcal{L}_0 \Rightarrow$ scalar d.o.f., $\mathcal{L}_{\frac{1}{2}} \Rightarrow$ fermion d.o.f., $\mathcal{L}_1 \Rightarrow$ vector d.o.f., $\mathcal{L}_m \Rightarrow$ fermion mass term, $\mathcal{L}_{mix} \Rightarrow$ mixing term, $\mathcal{L}_d \Rightarrow$ kinetic D term, $\mathcal{L}_{4f} \Rightarrow$ four

fermion term and the SUGRA potential in *Jordan frame* is given by [10, 11]

$$V_J = \frac{\Phi^2(z, \bar{z})}{9} \left[e^{\mathcal{K}(z, \bar{z})} \left\{ (\nabla_\alpha \mathcal{W}(z)) G^{\alpha\bar{\beta}} (\nabla_{\bar{\beta}} \bar{\mathcal{W}}(z)) - 3|\mathcal{W}(z)|^2 \right\} + \frac{1}{2} (\mathbf{Re} f(z))^{-1 AB} \mathcal{P}_A \mathcal{P}_B \right] \quad (2.3)$$

where $\alpha = 1, 2, \dots, n$ represent number of complex scalars in the SUGRA chiral multiplet, $\mathcal{K}(z, \bar{z})$ is the Kähler potential, $\mathcal{W}(z)$ is the holomorphic superpotential, $f_{AB}(z)$ is the holomorphic kinetic gauge matrix field and the Killing potential or momentum map is denoted by \mathcal{P}_A which includes all the Yang-Mills transformation of the scalars through which *Fayet-Iliopoulos* terms are also taken care of. In equation(2.1) the supergravity *verbien* (inverse of *fünfbien*) is characterized by the transformation rule [33]

$$g_{\mu\nu}^J := \eta_{\hat{A}\hat{B}} \left(V_\mu^{\hat{A}} \otimes V_\nu^{\hat{B}} \right) \quad (2.4)$$

with

$$\text{Det}(V) = \sqrt{-g_J} = e_{(4)}. \quad (2.5)$$

Here we use the following definition of covariant derivative:

$$\nabla_\alpha \mathcal{W} := W_\alpha + \mathcal{K}_\alpha \mathcal{W} \quad (2.6)$$

where the subscript α signifies differentiation with respect to complex field z^α . By setting $\Phi = -3$ in the SUGRA action in *Jordan frame* reduces to the well known action in the *Einstein frame*. Consequently the potential stated in equation(2.3) can be related to its *Einstein frame* counterpart as

$$V_J = \frac{\Phi^2(z, \bar{z})}{9} V_E, \quad (2.7)$$

where the J and E subscript are used for *Jordan* and *Einstein* frame. Here both the frames are connected via the *superconformal transformation* defined in terms of the metric as

$$g_{\mu\nu}^J = \Omega^2(z, \bar{z}) g_{\mu\nu}^E \quad (2.8)$$

where we identify the conformal factor with

$$\Omega^2(z, \bar{z}) = -\frac{\Phi(z, \bar{z})}{3} = e^{-\frac{\mathcal{K}(z, \bar{z})}{3}} \quad (2.9)$$

which yields a purely bosonic action in $\mathcal{N}=1, \mathcal{D}=4$ SUGRA in a specific *Jordan frame* triggering the *superHiggs mechanism*. The SUGRA action includes $SU(2, 2|1)$ superconformal symmetry, local dilation, special conformal symmetry, special SUSY and local $U(1)_R$ symmetry and other local symmetries of $\mathcal{N}=1, \mathcal{D}=4$ SUGRA. Such a superconformal mechanism is very useful to embed a class of scale invariant Global Supersymmetric (GSUSY) models into SUGRA theory. By “embedding”, here we actually point towards the fact that the $\mathcal{N}=1, \mathcal{D}=4$ self-interacting SUGRA multiplets has a local Poincare SUSY which can be obtained by the breakdown of above mentioned superconformal symmetry. Consequently the pure SUGRA sector in the action stated by equation(2.1) breaks superconformal symmetry and the matter part remains superconformal after gauge fixing. In general, that the kinetic term is non-canonical is guaranteed by the following choice of *superconformal factor* [10–12]:

$$\Omega^2(z, \bar{z}) = 1 - \frac{1}{3} \left(\delta_{\alpha\bar{\beta}} z^\alpha \bar{z}^{\bar{\beta}} + \mathcal{J}(z) + \bar{\mathcal{J}}(\bar{z}) \right) \quad (2.10)$$

where $\mathcal{J}(z)$ and $\bar{\mathcal{J}}(\bar{z})$ are the phenomenological holomorphic functions considered in the Kähler gauge. It is important to mention here that the dilation symmetry implies $\Omega^2(z, \bar{z})$ to be homogeneous of first degree in both z and \bar{z} , $\mathcal{W}(z)$ to be homogeneous of third degree in z . Additionally local $U(1)_R$ symmetry implies $\Omega^2(z, \bar{z})$ is neutral and $\mathcal{W}(z)$ has chiral weight three (which has been taken care of in equation(2.12)).

Now using equation(2.9) and equation(2.10) one can find out the explicit expressions for SUGRA *frame function* and the *Kähler potential* in this context. Using these results we get following expression for the *Kähler* metric:

$$G^{\alpha\bar{\beta}} = \left(\frac{\partial^2 \Omega^2(z, \bar{z})}{\partial z_\alpha \partial \bar{z}_\beta} \right) = \left\{ 1 - \frac{1}{3} \left(\delta_{\alpha\bar{\beta}} z^\alpha \bar{z}^{\bar{\beta}} + \mathcal{J}(z) + \bar{\mathcal{J}}(\bar{z}) \right) \right\} \left[\delta^{\alpha\bar{\beta}} - \frac{1}{3} \left(z^\alpha \bar{z}^{\bar{\beta}} + \delta^{\alpha\bar{\beta}} (\mathcal{J}(z) + \bar{\mathcal{J}}(\bar{z})) \right) \right] \quad (2.11)$$

Assuming the non-canonical structure of the *superconformal factor* stated in equation(2.10) let us now prove the equivalence of F-term SUGRA potential in superconformal *Jordan frame* and in GSUSY. We start with a renormalizable $\mathcal{N}=1$, $\mathcal{D}=4$ SUGRA where the most generalized expression of the superpotential is constrained to the following cubic form:

$$\mathcal{W}(z) = \frac{1}{3} \mathbf{d}_{\alpha\beta\gamma} z^\alpha z^\beta z^\gamma \quad (2.12)$$

where $\mathbf{d}_{\alpha\beta\gamma}$'s are the trilinear couplings in SUGRA theory. Equation(2.12) breaks the $\mathcal{SU}(1, \mathbf{n})$ symmetry. Now considering the fact that the SUGRA superpotential is homogeneous of the third degree in z^α 's we get:

$$\begin{aligned} \mathcal{W}_\alpha z^\alpha &= 3\mathcal{W}, \\ \bar{\mathcal{W}}_{\bar{\alpha}} \bar{z}^{\bar{\alpha}} &= 3\bar{\mathcal{W}}. \end{aligned} \quad (2.13)$$

Considering all the above facts the *Jordan frame* potential turns out to be

$$\begin{aligned} V_J^F &= \left(1 - \frac{1}{3} (\mathcal{J}(z) + \bar{\mathcal{J}}(\bar{z}))\right) [V_{GSUSY}^F + \bar{\mathcal{W}} (\partial_{z^\alpha} \mathcal{J}(z)) + \mathcal{W} (\partial_{\bar{z}^{\bar{\alpha}}} \bar{\mathcal{J}}(\bar{z}))] + |\mathcal{W}|^2 \left\{ \delta_{\alpha\bar{\beta}} z^\alpha \bar{z}^{\bar{\beta}} + \mathcal{J}(z) + \bar{\mathcal{J}}(\bar{z}) \right. \\ &\quad \left(1 - \frac{1}{3} (\mathcal{J}(z) + \bar{\mathcal{J}}(\bar{z}))\right) \left[\delta_{\bar{\gamma}\lambda} \bar{z}^{\bar{\gamma}} z^\lambda + z^\alpha (\partial_{z^\alpha} \mathcal{J}(z)) + \bar{z}^{\bar{\beta}} (\partial_{\bar{z}^{\bar{\alpha}}} \bar{\mathcal{J}}(\bar{z})) + \delta^{\alpha\bar{\beta}} (\partial_{z^\alpha} \mathcal{J}(z)) (\partial_{\bar{z}^{\bar{\alpha}}} \bar{\mathcal{J}}(\bar{z})) \right] \\ &\quad \left. - \frac{1}{3} z^\alpha \bar{z}^{\bar{\beta}} \left[\delta_{\alpha\bar{\gamma}} \delta_{\bar{\beta}\alpha'} \bar{z}^{\bar{\gamma}} z^{\alpha'} + \delta_{\bar{\beta}\alpha'} z^{\alpha'} (\partial_{z^\alpha} \mathcal{J}(z)) + \delta_{\alpha\bar{\gamma}} \bar{z}^{\bar{\gamma}} (\partial_{\bar{z}^{\bar{\beta}}} \bar{\mathcal{J}}(\bar{z})) + (\partial_{z^\alpha} \mathcal{J}(z)) (\partial_{\bar{z}^{\bar{\beta}}} \bar{\mathcal{J}}(\bar{z})) \right] \right\} \\ &\quad - \frac{1}{3} z^\alpha \bar{z}^{\bar{\beta}} \{ 3|\mathcal{W}|^2 \delta_{\alpha\bar{\beta}} + \mathcal{W} \bar{\mathcal{W}}_{\bar{\beta}} (\partial_{z^\alpha} \mathcal{J}(z)) + \delta_{\bar{\beta}\gamma} \bar{\mathcal{W}} \mathcal{W}_\alpha z^\gamma + \mathcal{W} \mathcal{W}_\alpha (\partial_{\bar{z}^{\bar{\beta}}} \bar{\mathcal{J}}(\bar{z})) \} \end{aligned} \quad (2.14)$$

where GSUSY potential $V_{GSUSY} = \delta^{\alpha\bar{\beta}} \mathcal{W}_\alpha \bar{\mathcal{W}}_{\bar{\beta}}$. Here the superscript F implies the F-term potential. Here it is important to mention that when *superconformal symmetry is gauge fixed*, the matter multiplets are preserved, which implies $\mathcal{J}(z) = 0$ and $\bar{\mathcal{J}}(\bar{z}) = 0$. Consequently equation(2.14) reduces to the following form:

$$V_J^F = V_{GSUSY}^F - \frac{1}{3} \delta_{\alpha\bar{\gamma}} \delta_{\bar{\beta}\alpha'} z^\alpha \bar{z}^{\bar{\beta}} \bar{z}^{\bar{\gamma}} z^{\alpha'} |\mathcal{W}|^2. \quad (2.15)$$

Now demanding renormalizability of the potential in SUGRA theory it is evident from equation(2.15) that the contribution from the last term is absent and we have:

$$V_J^F \simeq V_{GSUSY}^F \quad (2.16)$$

leading to the equivalence of F-term potentials as claimed above. Next we will concentrate on a specific situation where the superconformal symmetry is broken via the non-minimal coupling parameter χ with gravity. Consequently the frame function stated in equation(2.10) is modified as [10, 11]:

$$\Omega^2(z, \bar{z}) = -|z^0|^2 + |z^\alpha|^2 - \chi \left(\Theta_{\alpha\beta} \frac{z^\alpha z^\beta \bar{z}^{\bar{0}}}{z^0} + \bar{\Theta}_{\alpha\beta} \frac{\bar{z}^{\bar{\alpha}} \bar{z}^{\bar{\beta}} z^0}{\bar{z}^{\bar{0}}} \right) \quad (2.17)$$

which characterizes the non-flat moduli space geometry in SUGRA. Now gauge fixing criteria demands that in Planckian Unit system the compensator fields satisfy $z^0 = \bar{z}^{\bar{0}} = \sqrt{3}$. This implies a subsequent modification in the matter part of the inverse Kähler metric of the enlarged space can be expressed as:

$$\begin{aligned} G^{\alpha\bar{\beta}} &= \delta^{\alpha\bar{\beta}} - \frac{4\chi^2 \delta^{\alpha\bar{\lambda}} \delta^{\sigma\bar{\beta}} \Theta_{\sigma\zeta} \bar{\Theta}_{\bar{\lambda}\bar{\xi}} z^\zeta \bar{z}^{\bar{\xi}}}{[3 - \chi (\Theta_{\gamma\eta} z^\gamma z^\eta + \bar{\Theta}_{\bar{\gamma}\bar{\eta}} \bar{z}^{\bar{\gamma}} \bar{z}^{\bar{\eta}}) + 4\chi^2 \delta^{\gamma\bar{\eta}} \Theta_{\gamma\zeta} \bar{\Theta}_{\bar{\eta}\bar{\rho}} z^\rho \bar{z}^{\bar{\rho}}]}, \\ G^{0\bar{0}} &= - \frac{2\sqrt{3}\chi \delta^{\lambda\bar{\beta}} \Theta_{\lambda\xi} z^\xi}{[3 - \chi (\Theta_{\gamma\eta} z^\gamma z^\eta + \bar{\Theta}_{\bar{\gamma}\bar{\eta}} \bar{z}^{\bar{\gamma}} \bar{z}^{\bar{\eta}}) + 4\chi^2 \delta^{\gamma\bar{\eta}} \Theta_{\gamma\rho} \bar{\Theta}_{\bar{\eta}\bar{\sigma}} z^\rho \bar{z}^{\bar{\sigma}}]}, \\ G^{\alpha\bar{0}} &= - \frac{2\sqrt{3}\chi \delta^{\alpha\bar{\lambda}} \bar{\Theta}_{\bar{\lambda}\bar{\xi}} \bar{z}^{\bar{\xi}}}{[3 - \chi (\Theta_{\gamma\eta} z^\gamma z^\eta + \bar{\Theta}_{\bar{\gamma}\bar{\eta}} \bar{z}^{\bar{\gamma}} \bar{z}^{\bar{\eta}}) + 4\chi^2 \delta^{\gamma\bar{\eta}} \Theta_{\gamma\rho} \bar{\Theta}_{\bar{\eta}\bar{\sigma}} z^\rho \bar{z}^{\bar{\sigma}}]}, \\ G^{0\bar{\beta}} &= - \frac{3}{[3 - \chi (\Theta_{\gamma\eta} z^\gamma z^\eta + \bar{\Theta}_{\bar{\gamma}\bar{\eta}} \bar{z}^{\bar{\gamma}} \bar{z}^{\bar{\eta}}) + 4\chi^2 \delta^{\gamma\bar{\eta}} \Theta_{\gamma\zeta} \bar{\Theta}_{\bar{\eta}\bar{\rho}} z^\rho \bar{z}^{\bar{\rho}}]}, \end{aligned} \quad (2.18)$$

subject to the *orthonormalization condition*

$$G^{0\bar{\beta}} G_{0\bar{\gamma}} + G^{\alpha\bar{\beta}} G_{\alpha\bar{\gamma}} = \delta_{\bar{\gamma}}^{\bar{\beta}}. \quad (2.19)$$

This will directly modify the *Jordan frame* potential stated in equation(2.3). In the next two sections we will discuss elaborately the cosmological consequences of such non-minimal coupling parameter in the context of *superHiggs* theory.

III. INFLATIONARY MODEL BUILDING FOR DIFFERENT VALUES OF THE NON-MINIMAL COUPLING (χ)

In this section we will start our discussion with a simple gauge fixed version of frame function in presence of a superconformal symmetry breaking term (χ) in the Planckian unit:

$$\Phi(H, S, \bar{H}, \bar{S}) = -3 - \frac{1}{4} \left(1 + \frac{3\chi}{2}\right) [(H - \bar{H})^2 + (S - \bar{S})^2] + \frac{1}{4} \left(1 - \frac{3\chi}{2}\right) [(H + \bar{H})^2 + (S + \bar{S})^2]. \quad (3.1)$$

Using equation(2.9) the conformal factor turns out to be:

$$\Omega^2(H, S, \bar{H}, \bar{S}) = 1 + \frac{1}{12} \left(1 + \frac{3\chi}{2}\right) [(H - \bar{H})^2 + (S - \bar{S})^2] - \frac{1}{12} \left(1 - \frac{3\chi}{2}\right) [(H + \bar{H})^2 + (S + \bar{S})^2]. \quad (3.2)$$

Here the superHiggs sector $H = \frac{H_1 + iH_2}{\sqrt{2}}$ and $S = \frac{S_1 + iS_2}{\sqrt{2}}$ are complex scalar fields in the SUGRA chiral multiplet. Depending on the numerical values of χ , shift symmetry of H and S field is preserved. The behavior of the Higgs potential for various values of the non-minimal coupling is explicitly shown in figure(1). This shows that as the strength of the non-minimal coupling increases, the corresponding potential becomes more flat. In the next subsections we will study the cosmological consequences of these models in detail.

A. Models with $\chi = \frac{2}{3}$

In this branch the conformal factor is given by:

$$\Omega^2(H, S, \bar{H}, \bar{S}) = 1 + \frac{1}{6} [(H - \bar{H})^2 + (S - \bar{S})^2] \quad (3.3)$$

which is connected to the Kähler potential via equation(2.9). In this context the following transformations

$$\begin{aligned} H &\rightarrow H + C_H, \\ S &\rightarrow S + C_S \end{aligned} \quad (3.4)$$

lead to the shift symmetry of the Kähler potential with respect to $(H - \bar{H})$ and $(S - \bar{S})$, provided C_H and C_S are constant shifts along real axis of H and S complex plane.

Class of models	Ω^2	\mathcal{W}	V_J	V_E
H real, S=0	1	$-\lambda_1 S \left(H\bar{H} - \frac{v_1^2}{2}\right)$	$\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2$	$\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2$
H=0, S real	1	$-\lambda_2 H \left(S\bar{S} - \frac{v_2^2}{2}\right)$	$\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2$	$\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2$
H complex, S=0	$\left(1 - \frac{H_2^2}{3}\right)$	$-\lambda_1 S \left(H\bar{H} - \frac{v_1^2}{2}\right)$	$\frac{\lambda_1^2}{4} (H_1^2 + H_2^2 - v_1^2)^2$	$\frac{\lambda_1^2}{4} \frac{(H_1^2 + H_2^2 - v_1^2)^2}{\left(1 - \frac{H_2^2}{3}\right)^2}$
H=0, S complex	$\left(1 - \frac{S_2^2}{3}\right)$	$-\lambda_2 H \left(S\bar{S} - \frac{v_2^2}{2}\right)$	$\frac{\lambda_2^2}{4} (S_1^2 + S_2^2 - v_2^2)^2$	$\frac{\lambda_2^2}{4} \frac{(S_1^2 + S_2^2 - v_2^2)^2}{\left(1 - \frac{S_2^2}{3}\right)^2}$

TABLE I: Jordan frame and Einstein frame potentials obtained from $\chi = \frac{2}{3}$ branch.

In table(I) we have listed several class of *Jordan frame* and *Einstein frame* potentials from all possible physical combinations of H and S obtained from superconformal transformation mentioned in equation(3.5). In this article, potentials obtained from H and S in any branch are exactly similar. So we will restrict ourselves to the H dependent models for cosmological parameter estimation. To confront with the *recently observed Higgs at LHC here we fix the VEV, $v_1 = 246 \text{ GeV}$ with mass 125 GeV .*

B. Models with $\chi = -\frac{2}{3}$

In this branch the conformal factor in equation(3.2) reduces to the following:

$$\Omega^2(H, S, \bar{H}, \bar{S}) = 1 - \frac{1}{6} [(H + \bar{H})^2 + (S + \bar{S})^2] \quad (3.5)$$

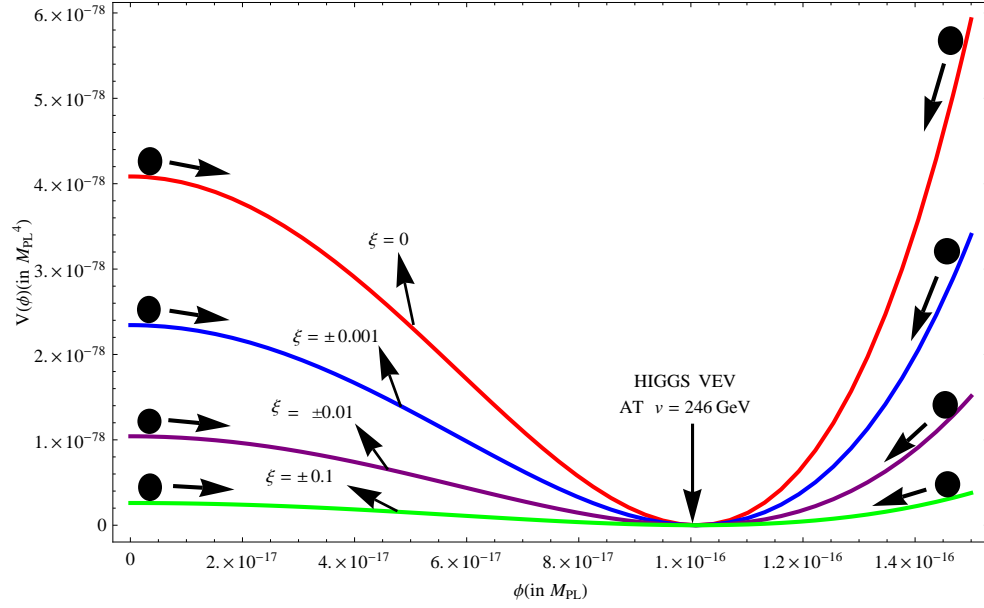


FIG. 1: Family of Higgs potentials for different numerical values of non-minimal coupling ξ , starting from $\xi = 0$. Inflation occurs either when the field $\phi = (H, S)$ rolls down from its large numerical values or when it rolls down from $\phi = 0$.

Potential	Confronts with	Coupling (λ_1) ($\times 10^{-7}$)	P_S ($\times 10^{-9}$)	n_S	α_S ($\times 10^{-4}$)	r	Ω_Λ	Ω_m	σ_8	η_{Rec} Mpc	η_0 Mpc
$\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2$	Λ CDM(WMAP9 + spt + act + snls3 + bao)	3.969	2.475	0.951	-7.931	0.260	0.711	0.289	0.807	283.03	14292.70
$\frac{\lambda_1^2}{4} \frac{(H_1^2 + H_2^2 - v_1^2)^2}{(1 - \frac{H_2^2}{3})^2}$	Λ CDM(WMAP9 + spt + act + snls3 + bao)	3.969	2.475	0.951	-7.931	0.260	0.711	0.289	0.807	283.15	14292.7

TABLE II: Cosmological parameter estimation for observationally allowed models obtained from $\chi = \frac{2}{3}$ branch.

which is connected to the Kähler potential via equation(2.9). In this context the following transformations

$$\begin{aligned} H &\rightarrow H + \tilde{C}_H, \\ S &\rightarrow S + \tilde{C}_S \end{aligned} \quad (3.6)$$

lead to the shift symmetry of the Kähler potential with respect to $(H + \tilde{H})$ and $(S + \tilde{S})$, provided \tilde{C}_H and \tilde{C}_S are constant shifts along imaginary axis of H and S complex plane.

Class of models	Ω^2	\mathcal{W}	V_J	V_E
H real, S=0	$(1 - \frac{H_1^2}{3})$	$-\lambda_1 S (H\bar{H} - \frac{v_1^2}{2})$	$\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2$	$\frac{\lambda_1^2}{4} \frac{(H_1^2 - v_1^2)^2}{(1 - \frac{H_1^2}{3})^2}$
H=0, S real	$(1 - \frac{S_1^2}{3})$	$-\lambda_2 H (S\bar{S} - \frac{v_2^2}{2})$	$\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2$	$\frac{\lambda_2^2}{4} \frac{(S_1^2 - v_2^2)^2}{(1 - \frac{S_1^2}{3})^2}$
H complex, S=0	$(1 - \frac{H_1^2}{3})$	$-\lambda_1 S (H\bar{H} - \frac{v_1^2}{2})$	$\frac{\lambda_1^2}{4} (H_1^2 + H_2^2 - v_1^2)^2$	$\frac{\lambda_1^2}{4} \frac{(H_1^2 + H_2^2 - v_1^2)^2}{(1 - \frac{H_1^2}{3})^2}$
H=0, S complex	$(1 - \frac{S_1^2}{3})$	$-\lambda_2 H (S\bar{S} - \frac{v_2^2}{2})$	$\frac{\lambda_2^2}{4} (S_1^2 + S_2^2 - v_2^2)^2$	$\frac{\lambda_2^2}{4} \frac{(S_1^2 + S_2^2 - v_2^2)^2}{(1 - \frac{S_1^2}{3})^2}$

TABLE III: Jordan frame and Einstein frame potentials obtained from $\chi = -\frac{2}{3}$ branch.

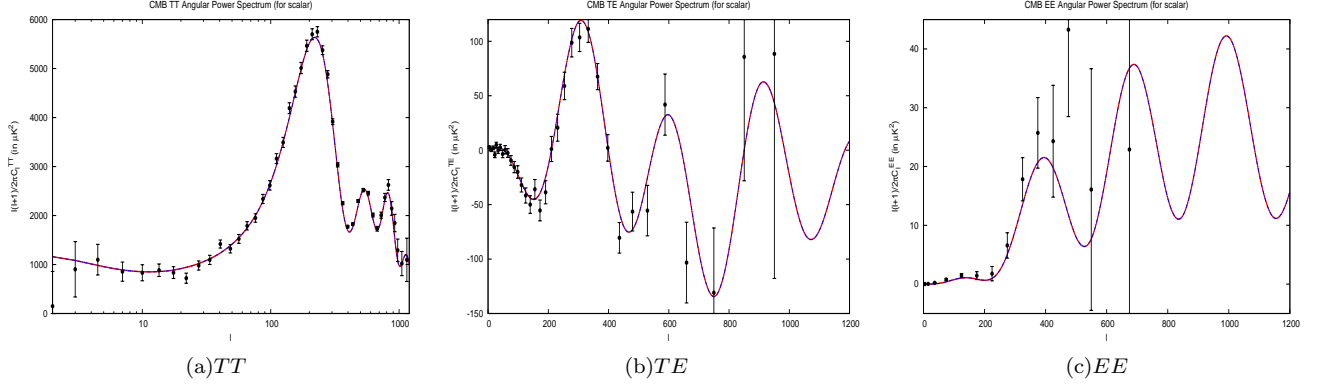


FIG. 2: Variation of CMB angular power spectra vs multipoles (l) for (a) TT, (b) TE and (c) EE mode from $\chi = \frac{2}{3}$ branch.

Potential	Confronts with	Coupling(λ_1) ($\times 10^{-7}$)	P_R ($\times 10^{-9}$)	n_s	α_s ($\times 10^{-4}$)	r	Ω_Λ	Ω_m	σ_8	η_{Rec} Mpc	η_0 Mpc
$\frac{\lambda_1^2}{4} \left(H_1^2 + H_2^2 - v_1^2 \right)^2$ $\left(1 - \frac{H_1^2}{3} \right)^2$	Λ CDM(WMAP9 + spt +act + snls3 + bao)	3.969	2.475	0.951	-7.932	0.260	0.711	0.289	0.807	283.15	14292.7

TABLE IV: Cosmological parameter estimation from observationally feasible model obtained from $\chi = -\frac{2}{3}$ branch.

In table(III) we mention the various class of *Jordan frame* and *Einstein frame* potentials from all possible physical combinations of H and S obtained from superconformal transformation mentioned in equation(3.5).

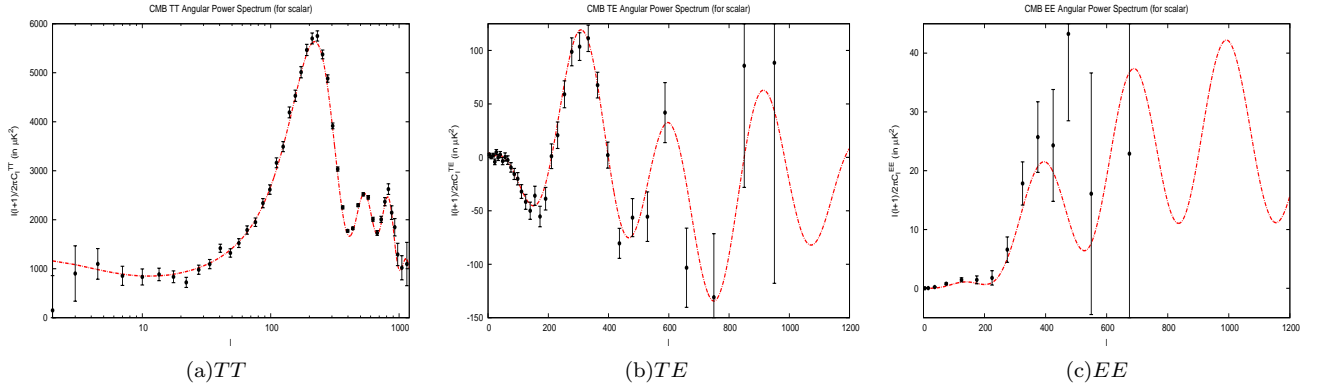


FIG. 3: Variation of CMB angular power spectra vs multipoles (l) for (a) TT, (b) TE and (c) EE mode from $\chi = -\frac{2}{3}$ branch.

In table(II) and table(IV) we mention all the cosmological parameters estimated from the observationally allowed potentials from $\chi = \frac{2}{3}$ and $\chi = -\frac{2}{3}$ branch respectively. From the numerical analysis we have explicitly shown that almost all of these proposed models confront well with latest *WMAP9* data combined with several complementary dataset obtained from *SPT*, *ACT*, *SNLS3*, *BAO* and h_0 observations in the Λ CDM or Λ CDM + *run* background. Next implementing the information obtained for each model from the cosmological code *CAMB* we estimate dark energy density (Ω_Λ), matter density (Ω_m) and its r.m.s. fluctuation (σ_8) etc. Hence we plot the behavior of CMB angular power spectrum for *TT*, *TE* and *EE* polarization obtained from $\chi = \pm \frac{2}{3}$ branch as shown in figure(2(a)-2(c)) and figure(3(a)-3(c)) for scalar mode.

In this article our prime objective is to study the cosmological consequences of single field inflationary potentials. For such cases the fields other than inflaton (i.e. background fields) can trigger the two phenomenological scenarios : preheating and reheating. Further this will directly or indirectly effect the leptogenesis [28, 29, 34–37] and baryogenesis [38, 39] scenario depending on the strength of the different decay channels of the inflatons into different particle constituents and the corresponding CP asymmetry of different branch. For the precise estimation of cosmological

parameters here we fix the value of all the background fields at GUT scale ($H_1 = H_2 = 0.9 \times 10^{16} \text{ GeV}$).

In this context, all the potentials are derived from SUGRA or from its superconformal extension. Consequently the energy scale of the potentials is around GUT scale. This directly satisfies the constraint on energy scale as $\mu_{GUT} < \Lambda_{UV}$, where $\mu_{GUT} \sim 10^{16} \text{ GeV}$ is the corresponding energy scale of SUGRA and $\Lambda_{UV} = M_{PL}$ be the UV(Ultra-Violet) cut-off theory. Here all the *Yukawa* type couplings (λ_1, λ_2) are energy scale dependent which will follow the *Renormalization Group* (RG) flow [40] via *Callan-Symanzik* equation. For the numerical estimation we fix the values of the *Yukawa* type couplings at GUT scale in the present context. Moreover, after applying RG flow from GUT to EWSB scale all of them becomes large ($\sim 2.065 \times 10^{-3}$) imposing the experimental constraints from *LHC*. It is a ray of hope for near future that proper bound on the self coupling is measurable in the next run of the *LHC*. Most importantly, the very recent Higgs mass bound observed at *LHC* and latest observational data from *WMAP9* have already ruled out the possibility of all of the proposed inflationary potentials at the EWSB scale in absence of any symmetry breaking non-minimal coupling. In this article by thorough numerical analysis we explicitly show that without introducing any non-minimal coupling all the proposed inflationary potentials obtained from the $\chi = \pm \frac{2}{3}$ branches are observationally favored at the GUT scale. On the other hand such running in the *Yukawa* type of couplings are inducing the possibility of *Primordial Black Hole* (PBH) formation [9, 41, 42] depending on the running on the model dependent cosmological parameter α_s . A very interesting fact for the inflationary model building is that the present observation from PLANCK (using WMAP9 data as a prior and the complementary data set (PLANCK lensing+CMB high l +BAO) [43] has predicted α_S and κ_S to be -0.013 ± 0.009 (although at 1.5σ) and $0.020^{+0.016}_{-0.015}$ respectively. Additionally for both $\chi = \pm \frac{2}{3}$ branches tensor to scalar ratio (r) are within the observational upper bound of *WMAP9*.

C. Models with $\chi \neq \pm \frac{2}{3}$

In this context the symmetry breaking parameter χ is connected with the non-minimal coupling ξ present as $\frac{\xi}{2}\phi^2 R$ in the action. To explore more features from this sector we consider two physical situations given by:

$$\chi - \frac{2}{3} = 4\xi_1 \quad (3.7)$$

$$\chi + \frac{2}{3} = 4\xi_2 \quad (3.8)$$

where ξ_1 and ξ_2 are the two non-minimal couplings approaching from $\frac{2}{3}$ and $-\frac{2}{3}$.

From equation(3.7) and equation(3.8) the superconformal factors can be expressed as:

$$\Omega_1^2(H, \bar{H}, S, \bar{S}) = 1 + \frac{1}{2} \left(\xi_1 + \frac{1}{3} \right) [(H - \bar{H})^2 + (S - \bar{S})^2] + \frac{\xi_1}{2} [(H + \bar{H})^2 + (S + \bar{S})^2] \quad (3.9)$$

$$\Omega_2^2(H, \bar{H}, S, \bar{S}) = 1 + \frac{\xi_2}{2} [(H - \bar{H})^2 + (S - \bar{S})^2] + \frac{1}{2} \left(\xi_2 - \frac{1}{3} \right) [(H + \bar{H})^2 + (S + \bar{S})^2] \quad (3.10)$$

In table(III) and table(IV) we mention all types of inflationary potentials in *Jordan frame* and *Einstein frame* as obtained from the two possible physical branches of the superconformal transformations mentioned in equation(3.9) and equation(3.10).

Next we have mentioned all the cosmological parameters estimated from the $\chi \neq \frac{2}{3}$ ($\chi - \frac{2}{3} = 4\xi_1$ and $\chi + \frac{2}{3} = 4\xi_2$) branch in table(VI) and table(VIII). This clearly shows non-minimal coupling (ξ_1, ξ_2) dependent models confront with latest *WMAP9* data combined with several complementary datasets in the Λ CDM or Λ CDM + *run* background. We also show that if we allow the above mentioned non-minimal couplings along with very recent *LHC* Higgs mass bound and latest observational constraints from *WMAP9* then almost all of the proposed inflationary potentials are favored starting from EWSB to GUT scale depending on the RG flow in *Yukawa* type coupling. Throughout the numerical analysis we allow both the signatures of the non-minimal coupling. We also avoid specific values of the non-minimal couplings for which divergences are appearing in the proposed potentials. During the analysis we observe that only for $(\xi_1, \xi_2) > 0$ the first two models appearing in table(VIII) and table(VI) are in good agreement with latest observation. On the contrary for $(\xi_1, \xi_2) < 0$ only the third model confronts *WMAP9*. Moreover, for the numerical estimations we consider only those values of the non-minimal couplings for which the proposed models are free from any poles. The behavior of tensor to scalar ratio (r) with respect to the scalar spectral index (n_S) for all class of proposed models of inflation are depicted in figure(6).

Class of models	Ω_1^2	\mathcal{W}	V_J	V_E
H real, S=0	$(1 + \xi_1 H_1^2)$	$-\lambda_1 S \left(H\bar{H} - \frac{v_1^2}{2} \right)$	$\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2$	$\frac{\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2}{(1 + \xi_1 H_1^2)^2}$
H=0, S real	$(1 + \xi_1 S_1^2)$	$-\lambda_2 H \left(S\bar{S} - \frac{v_2^2}{2} \right)$	$\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2$	$\frac{\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2}{(1 + \xi_1 S_1^2)^2}$
H complex, S=0	$1 - (\xi_1 + \frac{1}{3}) H_2^2 + \xi_1 H_1^2$	$-\lambda_1 S \left(H\bar{H} - \frac{v_1^2}{2} \right)$	$\frac{\lambda_1^2}{4} (H_1^2 + H_2^2 - v_1^2)^2$	$\frac{\frac{\lambda_1^2}{4} (H_1^2 + H_2^2 - v_1^2)^2}{[1 - (\xi_1 + \frac{1}{3}) H_2^2 + \xi_1 H_1^2]^2}$
H=0, S complex	$1 - (\xi_1 + \frac{1}{3}) S_2^2 + \xi_1 S_1^2$	$-\lambda_2 H \left(S\bar{S} - \frac{v_2^2}{2} \right)$	$\frac{\lambda_2^2}{4} (S_1^2 + S_2^2 - v_2^2)^2$	$\frac{\frac{\lambda_2^2}{4} (S_1^2 + S_2^2 - v_2^2)^2}{[1 - (\xi_1 + \frac{1}{3}) S_2^2 + \xi_1 S_1^2]^2}$

TABLE V: *Jordan frame and Einstein frame potentials obtained from $\chi - \frac{2}{3} = 4\xi_1$ branch.*

Potential	Confronts with	Couplings ($\times 10^{-7}$)	ξ_1	P_R ($\times 10^{-9}$)	n_s	α_s ($\times 10^{-4}$)	r	Ω_Λ	Ω_m	σ_8	η_{Rec} Mpc	η_0 Mpc
$\frac{\lambda_1^2}{4} \frac{(H_1^2 - v_1^2)^2}{(1 + \xi_1 H_1^2)^2}$	Λ CDM(WMAP9 + spt + act + snls3 + h_0)	51.06	0.1	2.423	0.973	-4.700	0.015	0.745	0.255	0.812	286.26	14403.2
$\frac{\lambda_1^2}{4} \frac{(H_1^2 + H_2^2 - v_1^2)^2}{[1 - (\xi_1 + \frac{1}{3}) H_2^2 + \xi_1 H_1^2]^2}$	Λ CDM(WMAP9 + spt + act + snls3 + h_0)	51.06	0.1	2.423	0.973	-4.700	0.015	0.745	0.255	0.812	286.26	14403.2
$\frac{\lambda_2^2}{4} \frac{(H_1^2 + H_2^2 - v_2^2)^2}{[1 - (\xi_1 + \frac{1}{3}) H_2^2 + \xi_1 H_1^2]^2}$	Λ CDM(WMAP9 + spt + act + snls3 + h_0)	72.44	-0.5	2.422	0.973	-4.575	0.011	0.745	0.255	0.813	286.26	14403.2

TABLE VI: *Cosmological parameter estimation from observationally allowed models obtained from $\chi - \frac{2}{3} = 4\xi_1$ branch.*

IV. SUMMARY AND OUTLOOK

In this article we have proposed a class of supergravity motivated models to implement Higgs inflation, where the Higgs field is non-minimally coupled to the gravity sector via symmetry breaking coupling (χ). We have followed the analysis by making use of superconformal techniques in the Kähler manifold. Using such tools we have introduced a phenomenological Kähler potential which preserves shift symmetry for two minimal coupling $\chi = \pm \frac{2}{3}$ with gravity. This results in various classes of inflationary models which are made up of shift symmetry protected flat directions. We have elaborately discussed the consequences of superconformal techniques in the two preferred frame of references namely, Jordan and Einstein frames. Then we have explored the features of non-minimal coupling (ξ_1, ξ_2) connected with shift symmetry breaking branch $\chi \neq \frac{2}{3}$ in the context of Higgs inflation. Hence we have studied inflation from these proposed models by estimating the observable parameters originated from primordial quantum fluctuation for

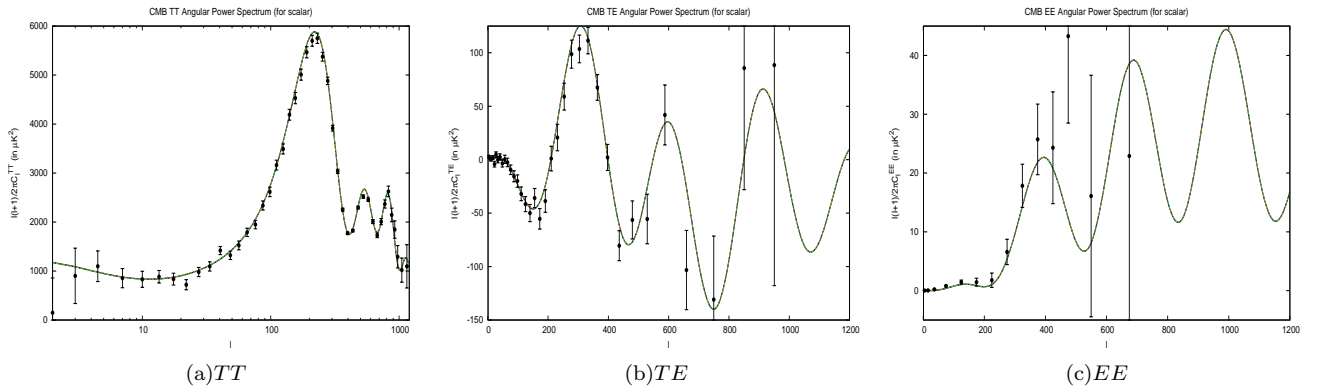


FIG. 4: Variation of CMB angular power spectra vs multipoles (l) (a) TT, (b) TE and (c) EE mode from $\chi - \frac{2}{3} = 4\xi_1$ branch.

Class of models	Ω_2^2	\mathcal{W}	V_J	V_E
H real, S=0	$[1 + (\xi_2 - \frac{1}{3}) H_1^2]$	$-\lambda_1 S \left(H\bar{H} - \frac{v_1^2}{2} \right)$	$\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2$	$\frac{\frac{\lambda_1^2}{4} (H_1^2 - v_1^2)^2}{[1 + (\xi_2 - \frac{1}{3}) H_1^2]^2}$
H=0, S real	$[1 + (\xi_2 - \frac{1}{3}) S_1^2]$	$-\lambda_2 H \left(S\bar{S} - \frac{v_2^2}{2} \right)$	$\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2$	$\frac{\frac{\lambda_2^2}{4} (S_1^2 - v_2^2)^2}{[1 + (\xi_2 - \frac{1}{3}) S_1^2]^2}$
H complex, S=0	$1 + (\xi_2 - \frac{1}{3}) H_1^2 - \xi_2 H_2^2$	$-\lambda_1 S \left(H\bar{H} - \frac{v_1^2}{2} \right)$	$\frac{\lambda_1^2}{4} (H_1^2 + H_2^2 - v_1^2)^2$	$\frac{\frac{\lambda_1^2}{4} (H_1^2 + H_2^2 - v_1^2)^2}{[1 + (\xi_2 - \frac{1}{3}) H_1^2 - \xi_2 H_2^2]^2}$
H=0, S complex	$1 + (\xi_2 - \frac{1}{3}) S_1^2 - \xi_2 S_2^2$	$-\lambda_2 H \left(S\bar{S} - \frac{v_2^2}{2} \right)$	$\frac{\lambda_2^2}{4} (S_1^2 + S_2^2 - v_2^2)^2$	$\frac{\frac{\lambda_2^2}{4} (S_1^2 + S_2^2 - v_2^2)^2}{[1 + (\xi_2 - \frac{1}{3}) S_1^2 - \xi_2 S_2^2]^2}$

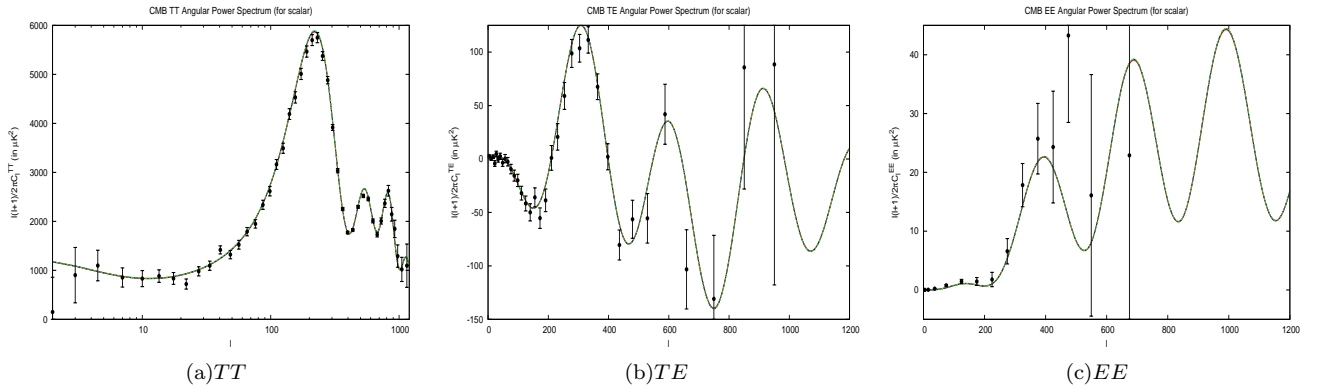
TABLE VII: Jordan frame and Einstein frame potentials obtained from $\chi + \frac{2}{3} = 4\xi_2$ branch.

Potential	Confronts with	Couplings ($\times 10^{-7}$)	ξ_2	F_R ($\times 10^{-9}$)	n_s	α_s ($\times 10^{-4}$)	r	Ω_Λ	Ω_m	σ_8	η_{Rec} Mpc	η_0 Mpc
$\frac{\lambda_1^2}{4} \frac{(H_1^2 - v_1^2)^2}{[1 + (\xi_2 - \frac{1}{3}) H_1^2]^2}$	Λ CDM(WMAP9 + spt + act + snls3 + h_0)	38.916	0.5	2.422	0.972	-4.822	0.018	0.745	0.255	0.811	286.26	14403.2
$\frac{\lambda_1^2}{4} \frac{(H_1^2 + H_2^2 - v_1^2)^2}{[1 + (\xi_2 - \frac{1}{3}) H_1^2 - \xi_2 H_2^2]^2}$	Λ CDM(WMAP9 + spt + act + snls3 + h_0)	72.42	0.5	2.421	0.973	-4.575	0.011	0.745	0.255	0.813	286.26	14403.2
$\frac{\lambda_2^2}{4} \frac{(H_1^2 + H_2^2 - v_2^2)^2}{[1 + (\xi_2 - \frac{1}{3}) H_1^2 - \xi_2 H_2^2]^2}$	Λ CDM(WMAP9 + spt + act + snls3 + h_0)	51.100	-0.1	2.427	0.973	-4.700	0.015	0.745	0.255	0.813	286.26	14403.2

TABLE VIII: Cosmological parameter estimation for observationally favored models obtained from $\chi + \frac{2}{3} = 4\xi_2$ branch.

scalar and tensor modes. We have further confronted our results with WMAP9 and various complementary dataset (*SPT, ACT, BAO, SNLS3, h_0*) by using CAMB. Further we have compared the behavior of theoretical CMB polarization power spectra for *TT, TE* and *EE* mode obtained from all of these proposed models with observational power spectra. We have also commented on the allowed range for the non-minimal couplings (ξ_1, ξ_2) and phenomenological *Yukawa* type of couplings which are very crucial inputs in the context of inflationary model building. This, collectively, provides an exhaustive study of the class of Higgs inflation from Kähler potential and consequently, their pros and cons.

An interesting open issue in this context is to study the role of Hiesenberg symmetry [44–46] in the present context. Other open issues is to study primordial black hole formation and its cosmological consequences from the running of the spectral index (α_S) and its running (κ_S) as the very recent *PLANCK* data gives an estimation at 1.5σ [43]. Moreover, the phenomenological consequences of all of these proposed models via reheating and leptogenesis are also a promising issue for future study.

FIG. 5: Variation of CMB angular power spectra vs multipoles (l) for (a) TT, (b) TE (c) EE mode from $\chi + \frac{2}{3} = 4\xi_2$ branch.

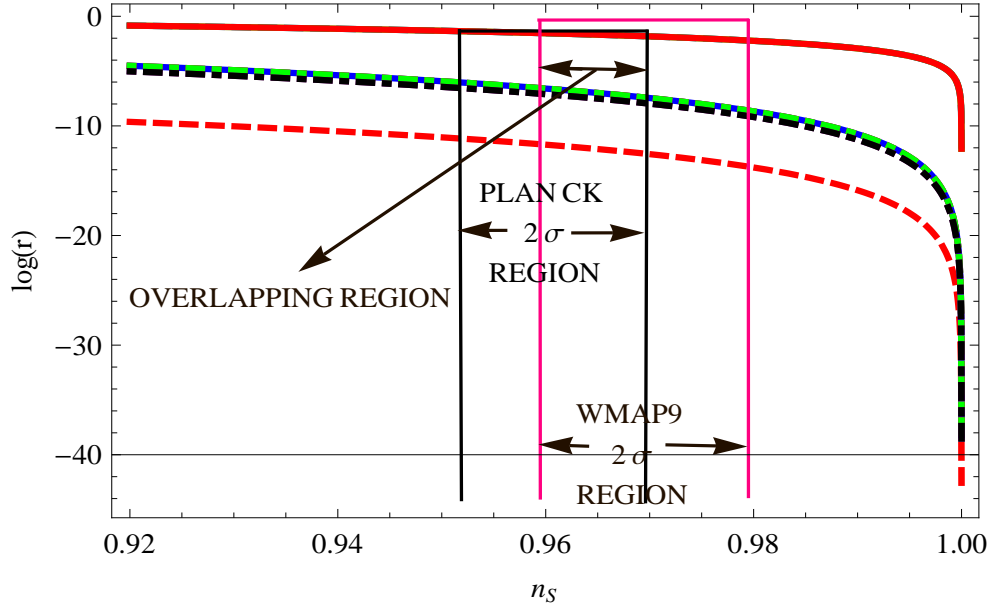


FIG. 6: Variation of tensor to scalar ratio (r) vs scalar spectral index (n_s) for the family of Higgs potentials for different numerical values of non-minimal coupling ξ . The value of the non-minimal coupling increases as we go down towards the plot. This also shows $\chi \pm \frac{2}{3} = 4\xi$ branches are more observationally favored compared to the $\chi = \pm \frac{2}{3}$ branches.

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